

Particle production with L-R neutrino oscillation

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When the Higgs field starts oscillation after Higgs inflation, gauge bosons are produced non-perturbatively near the Enhanced Symmetry Point (ESP). Just after the particle production, when the Higgs field is going away from the ESP, these gauge bosons gain mass and decay or annihilate into Standard Model (SM) fermions. Left-handed neutrinos can be generated in that way. If one assumes the see-saw mechanism, the mass matrix of a pair of left and right-handed neutrinos is non-diagonal. Although their mixing in the mass eigenstates is negligible in the true vacuum, it could be significant near the edge of the Higgs oscillation, where the off-diagonal component is large. Therefore, the left-handed neutrinos generated from the gauge bosons can start neutrino oscillation between the right-handed neutrinos. We study the particle production when such L-R neutrino oscillation is significant. For a working example, the non-thermal leptogenesis scenario after Higgs inflation is examined, which cannot be realized without the L-R neutrino oscillation. The same mechanism could be applied to other singlet particles whose abundance has been neglected.

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I. INTRODUCTION AND THE MODEL

The recent release of the Planck data [1] suggests that the Higgs inflation model [2, 3] is in agreement with the Planck constraints.¹ The idea of Higgs inflation is quite attractive since it uses the well-known scalar field and explains inflation without introducing a new field to the SM. On the other hand, it is hard to explain the baryon asymmetry of the current Universe within the SM. The neutrino mass, which is discovered to be non-zero in 1998 [6], is also beyond the SM. One way to explain the neutrino mass is to consider see-saw mechanism [7], which is also expected to explain the baryon asymmetry via leptogenesis [8]. In this paper we study a specific process of generating right-handed neutrinos during non-perturbative particle production (preheating) that may occur before reheating. We apply the idea of neutrino oscillation to the mass matrix of the see-saw mechanism. This idea works when the Higgs field has large vacuum condensate. Showing that significant mixing is possible during preheating, we demonstrate how the L-R neutrino oscillation works to generate right-handed neutrinos. Remember that conventional neutrino oscillation in the dense matter [9, 10] leads to a similar situation. In that case the matrix is given for a pair of ν_e and ν_μ . The diagonal element of ν_e becomes time-dependent due to the interaction between matter, while the other elements remain constant. In the core of the star, where the density is high, the mass eigenstates could be $\psi_+ \simeq \nu_e$ (heavy

state) and $\psi_- \simeq \nu_\mu$ (light state), while outside the star (in the vacuum) these relations turn out to be opposite. If $\nu_e \simeq \psi_+$ is produced in the core and it propagates adiabatically until it reaches the surface of the star, one will find emission of the heavier mass eigenstate ψ_+ from the star, which is now $\psi_+ \simeq \nu_\mu$. In our case, the left-handed neutrino (ν_L) is generated near the edge, where ν_L is a mixed state of ψ_+ and ψ_- . Of course, one will find $\psi_+ \simeq N_R$ in the true vacuum. Although our situation could be unusual, the use of the time-dependent matrix for the neutrino production is quite common in astro-particle physics. We apply our idea to the simplest leptogenesis scenario and find that the new mechanism assists the non-thermal leptogenesis scenario.

A. Preheating and neutrino oscillation after Higgs inflation

Since in this paper we have no space to review the details of the oscillation and the particle production after Higgs inflation, we carefully follow Ref.[3, 11] to avoid confusions. Here we prepared our starting point as simple as possible. We start with the effective potential

$$V(\chi) = \frac{1}{2}M^2\chi^2, \quad (1)$$

where $M \equiv \sqrt{\lambda}M_p/\sqrt{3}\xi$ and χ comes from the original Higgs field after conformal transformation. Here λ denotes the original quartic coupling of the Higgs field. The conformal factor of the transformation can be expressed using

$$\Omega^2(h) = 1 + \frac{\xi h^2}{M_p^2} \quad (2)$$

¹ The scale-dependence of the spectrum can be shifted by an additional degree of freedom [4]. Observation of the tensor mode and its spectral index (if possible) could fix such ambiguity [5].

$$\Omega^2(\chi) = e^{\alpha\kappa\chi}. \quad (3)$$

We used $\alpha \equiv \sqrt{2/3}$ and $\kappa \equiv M_p^{-1}$, where M_p is the reduced Planck mass of the original action. Here h denotes the condensate of the original Higgs field in the unitary gauge. The oscillation starts when $X_0 \sim 0.1M_p$, where X_n denotes the amplitude of the n -th oscillation. Consider the masses of “gauge bosons” in the Einstein frame;

$$\mathcal{L}_{W,Z}^E \sim \tilde{m}_W^2 \tilde{W}_\mu^+ \tilde{W}^{\mu-} + \frac{1}{2} \tilde{m}_Z^2 \tilde{Z}_\mu \tilde{Z}^\mu, \quad (4)$$

where

$$\tilde{m}_W^2 \equiv \frac{g_2^2 M_p^2 (1 - e^{-\alpha\kappa|\chi|})}{4\xi} \quad (5)$$

$$\tilde{m}_Z^2 \equiv \frac{\tilde{m}_W^2}{\cos\theta_W}. \quad (6)$$

Here g_2 is the $SU(2)_L$ gauge coupling constant at the corresponding scale and θ_W is the Weinberg angle defined at the scale. Tilde denotes quantities in the Einstein frame.

Similarly, the interaction between fermions and gauge bosons can be expressed as

$$\sim \frac{g_2}{\sqrt{2}} \tilde{W}_\mu^+ \tilde{J}^{\mu-} + \frac{g_2}{\sqrt{2}} \tilde{W}_\mu^- \tilde{J}^{\mu+} + \frac{g_2}{\cos\theta_W} \tilde{Z}_\mu \tilde{J}^{\mu 0}, \quad (7)$$

where $\tilde{J}^{\mu(\pm,0)}$ are charged and neutral currents of the SM fermions, which may contain left-handed neutrinos.

For a given family of the quark sector the mass of the fermions (f : the family index is omitted) is found from the Yukawa coupling. In the Einstein frame it can be written in a schematic form $\tilde{m}_f \tilde{f} \tilde{f}$, where

$$\tilde{m}_f \equiv \frac{y_f M_p}{\sqrt{2\xi}} \left(1 - e^{-\alpha\kappa|\chi|}\right)^{1/2}. \quad (8)$$

On the other hand, the neutrino mass is given by the see-saw mechanism. We choose the mass matrix for the left (ν_L) and the right (N_R) handed neutrinos;

$$\mathcal{M} = \begin{bmatrix} 0 & m_D \\ m_D & M_R \end{bmatrix}, \quad (9)$$

where we take $m_D = y h$.² \mathcal{M} can be diagonalized using the rotation matrix $U(\theta)$ as $\mathcal{M}_{\text{diag}} = U^T \mathcal{M} U \equiv \text{diag}(M_-, M_+)$ to give³

$$\frac{1}{2} \begin{pmatrix} M_R - \sqrt{M_R^2 + 4m_D^2} & 0 \\ 0 & M_R + \sqrt{M_R^2 + 4m_D^2} \end{pmatrix}, \quad (10)$$

² Here we consider only the first generation for sake of simplicity. In reality the flavor structure is very important. It is well known that in thermal leptogenesis the evolution of the lepton asymmetry depends on the flavor structure. We need to introduce more than two right-handed neutrinos to explain the observed active neutrino mass differences. Moreover, CP violation requires at least three generation.

³ See also Ref.[21].

where the angle of the rotation is given by

$$\theta = \frac{1}{2} \tan^{-1} \left(\frac{2m_D}{M_R} \right). \quad (11)$$

Conformal transformation leads to

$$\tilde{M}_R^2 = M_R^2 \Omega^{-2} = M_R^2 e^{-\alpha\kappa|\chi|} \quad (12)$$

$$\tilde{m}_D^2 = \frac{y_\nu^2 M_p^2 (1 - e^{-\alpha\kappa|\chi|})}{\xi}. \quad (13)$$

Practically, these variables are simplified using $1 - e^{-\alpha\kappa|\chi|} \simeq \alpha\kappa|\chi|$.

We consider the simplest see-saw mechanism in which heavy N_R ($M_R \gg h$ in the vacuum) is responsible for the SM neutrino mass.

In terms of the mass eigenstates (ψ_\mp), left and right-handed neutrinos are written as

$$\nu_L = \psi_- \cos\theta - \psi_+ \sin\theta \quad (14)$$

$$N_R = \psi_- \sin\theta + \psi_+ \cos\theta. \quad (15)$$

Normally (in the vacuum) one will find $M_R \gg m_D$ that leads to $\theta \simeq m_D/M_R \simeq 0$. Therefore there is no significant left-right mixing in the see-saw sector of the vacuum state. On the other hand, although temporarily the mixing could be significant when m_D is as large as M_R .

In Ref.[12], one can find extensive study of the non-perturbative particle production due to the time-dependent non-diagonal mass matrix (i.e, when $\dot{U} \neq 0$). Ref.[13] studies a model in which the mass term is not time-dependent but the interaction between other fields causes particle production. Alternatively one may consider higher dimensional interaction [3, 14, 15]. However, direct production of fermions is not important in our case since fermions do not lead to parametric resonance.⁴

Besides neutrinos, other particles such as gauge bosons and quarks can be generated non-perturbatively. Among those, the most significant effect is expected for the gauge bosons because of the parametric resonance. However, since the gauge bosons may have huge mass when the amplitude of the oscillation is large, they can partially decay into fermions before the next particle production. This gives suppression of the resonance in the first stage. Then the gauge bosons accumulate slowly and finally the parametric resonance takes place. According to Ref.[3], the main process responsible for the energy transfer is the annihilation of the W bosons. Then the radiation starts to dominate when the amplitude is $X_r \sim M$. Thermalization takes place later at $X_R \leq X_r$.

Since the currents that couple to those gauge bosons contain neutrinos, SM neutrinos can be generated when gauge bosons decay or annihilate. **One thing that is**

⁴ Notice also that conventional fermion preheating [16] uses $m_f = M_0 + g\phi$ that vanishes at $\phi = -M_0/g$, while the mass of ψ_+ cannot vanish in the see-saw mechanism.

unusual here is that the ν_L -state, which is produced from the gauge bosons, can be a mixed state of the mass eigenstates ψ_{\pm} if the transfer proceeds near the edge.⁵ The condition $\tilde{m}_D > \tilde{M}_R$ can be satisfied until the amplitude decreases to reach $X_e \simeq \xi M_R^2/(y_\nu^2 M_p) \sim M_R^2/(y_\nu^2 M)$. Using $m_\nu \equiv y_\nu^2 v_{EW}^2/M_R$, we can write

$$X_e \simeq 10^{14} \left(\frac{M_R}{M} \right) \left(\frac{v_{EW}}{10^2 \text{GeV}} \right)^2 \left(\frac{0.1 \text{eV}}{m_\nu} \right) \text{GeV}. \quad (16)$$

Here we can assume that neutrinos generated from the gauge bosons are relativistic. (See Ref.[3].) Therefore, typical time scale of the neutrino oscillation is measured by the square mass difference between ψ_+ and ψ_- .⁶ This gives the simple estimation of the typical time scale $\sim \frac{\Delta M_{\pm}^2}{4E} \sim \frac{M_R \sqrt{M_R^2 + 4m_D^2}}{m_D} \sim M_R$, which is identical to the right-handed neutrino mass.

The most obvious realization of the scenario is to take $M_R > M$. In that case the neutrino oscillation is fast enough to make oscillation before the Higgs field goes back to the origin. If one compares our result with Ref.[3], one will find that tiny Yukawa couplings are not important in our case, but small M_R makes the oscillation length large and may prevent L-R oscillations. In this paper we focus on this simplest scenario and examine the non-thermal leptogenesis scenario.

Although we are not discussing the possibility in this paper, less obvious scenario could be possible when $M_R \sim H$, where H is the Hubble parameter during oscillation. In this case, the ‘‘amplitude’’ of the Higgs oscillation is large enough to keep the maximum mixing, and the time scale of the neutrino oscillation is comparable to the time scale of the cosmological evolution of the amplitude. Indeed, coarse graining the Higgs oscillation and consider the amplitude of the oscillation as the time-dependent parameter of the model, one could be able to find significant production of right-handed neutrinos when $M_R > H$. Since in this case the typical time scale of the L-R neutrino oscillation is not very short, one has to consider the quantum Zeno effect. In the hot and dense environment (when the scattering of SM neutrinos is significant) there is a competition between the oscillation length and the mean free path. If the oscillation length is much larger than the mean free path, the L-R transition probability is hindered. This is the quantum Zeno effect. Usually one has to introduce a damping term in the neutrino oscillation equations [22]. However,

in the complicated far-from equilibrium environment of preheating it is quite difficult to quantify the damping.⁷ We are going to reserve such issue for future study.

Besides that, it should be noted that partial translation could be important even if the neutrino oscillation is blocked halfway, since the L-R translation is much more efficient compared with the usual process. Importantly, the same mechanism can be applied to other singlet particles whose abundance has been neglected.

Eventually, gauge bosons can produce N_R via the unconventional ‘‘neutrino oscillation’’ and the production mechanism can affect cosmological scenarios after Higgs inflation.

B. Non-thermal Leptogenesis from the L-R neutrino oscillation

The main purpose of this paper is to present a novel mechanism of generating particles from a time-dependent off-diagonal element in the mass matrix. The mechanism will work in general situation. Indeed, the mechanism may work in Grand Unified Theory (GUT), which will have many non-diagonal mass matrices. The calculation is quite complicated and the study requires careful numerical calculation.

Nevertheless, we believe that our attempt to show an application of the mechanism in the simplest set-up is very useful, as far as it gives an intuitive estimation of the quantities under some reasonable assumptions. Among possible applications of the model, dark matter production and leptogenesis would be the most important. In this paper, we choose the simplest model of non-thermal leptogenesis in which the non-equilibrium decay of the right-handed neutrino generates baryon (lepton) asymmetry of the Universe. In our model, the right-handed neutrino is never thermalized and there is no wash-out of the asymmetry. This is our simplest set-up of the model. The leptogenesis proceeds when ψ_+ decays into χ and ψ_- , where ψ_+ is a mixed state of N_R and ν_L . ψ_- turns into ν_L in the vacuum.

According to Ref.[3, 11], χ starts oscillation at $X_0 \sim 0.1 M_p$ and parametric resonance is (partially) suppressed until $X_{PR} \sim 10^{15} \text{GeV}$. At X_{PR} , we expect that 1/10 of the inflaton energy is translated into gauge bosons. This estimation is taken from Ref. [11]. Therefore we assume

$$\rho_{Bs}(t_{PR}) \simeq 0.1 \rho_\chi(t_{PR}). \quad (17)$$

Then we assume that the energy transfer is accomplished and thermalized when $X_R \sim X_r \sim M$. This gives the reheating temperature $T_R \sim 10^{13} \text{GeV}$.⁸ During the

⁵ Consider expansion of h around the expectation value of the condensate; $h \rightarrow v + \hat{h}$. Expanding the original interaction ($yhN_R\nu_L$) using mass eigenstates (ψ_{\pm}), one will find that the coefficient of the interaction $\hat{h}\psi_+\psi_-$ is proportional to $\cos 2\theta$. On the other hand, if ψ_+ does not decay during the oscillation, it can be translated into N_R (adiabatic process). See also Ref.[9, 10].

⁶ See Eq.(10).

⁷ We are focusing on changes in the initial condition. Therefore, active-sterile neutrino oscillation after reheating is not considered here.

⁸ As is noted in Ref.[11], estimation of the reheating temperature is not simple. Here we took T_R from Ref.[3].

process, certain fraction of gauge bosons is continuously converted into fermions. In reality the fraction is time-dependent and the calculation of the number densities is highly non-local (i.e, they must be calculated as the time integral of a complex system). To avoid further complexity, we assume $X_{PR} \simeq X_e$ and define a new parameter $\epsilon_+ < 1$, assuming that the number density of ψ_+ just after $X_{PR} \simeq X_e$ is given by

$$n_+ \equiv \epsilon_+ n_{Bs}, \quad (18)$$

where $n_{Bs} \equiv n_Z + n_W$ is the number density of the gauge bosons. Since ψ_+ can decay into the Higgs and ψ_- , we can define ϵ_{CP} to estimate the produced lepton number

$$n_L \sim \epsilon_{CP} n_+ \quad (19)$$

$$\sim \epsilon_{CP} \epsilon_+ \frac{0.1 M^2 X_{PR}^2}{g_2 \sqrt{M X_{PR}}} \quad (20)$$

$$\sim \epsilon_{CP} \epsilon_+ (M X_{PR})^{3/2}, \quad (21)$$

where $m_Z \sim m_W \sim g_2 \sqrt{M X}$ is assumed for simplicity. In the vacuum, when the background is static, ϵ_{CP} is identical to the conventional leptogenesis scenario, while in the above case one has to calculate the integral of the system to obtain ϵ_{CP} . In that sense both ϵ_+ and ϵ_{CP} are not local. Actual computation of those parameters is very difficult, and it depends on the model parameters of the lepton sector.⁹ Nevertheless, speculation would be possible for these parameters, which will be helpful for our purpose in this section. Since the decay or the annihilation before t_{PR} is sufficient to prevent the resonance, $\epsilon_+ \gtrsim 0.01 - 0.001$ would be conceivable.¹⁰ If the conventional CP violation is valid for our case, one can expect $\epsilon_{CP} \sim 10^{-5}$, which can be enhanced. Here the family multiplicity is implicitly assumed for N_R . If we can assume $n_L \propto a(t)^{-3}$ and $\rho_\chi \propto a(t)^{-3}$ between X_e and X_r , where $a(t)$ is the scale factor of the Universe, we find using $n_L(t_{PR}) = \epsilon_{CP} \epsilon_+ n_{Bs} = \epsilon_{CP} \epsilon_+ \rho_B / m_B = 0.1 \times \epsilon_{CP} \epsilon_+ \rho_\chi / m_B$ and $\rho_\chi(T_r) \simeq \rho_\chi(T_R) \simeq 10^2 T_R^4$:

$$\frac{n_L(t_R)}{T_R^3} \sim 0.1 \times \epsilon_{CP} \epsilon_+ \frac{n_L(t_{PR})}{n_L(t_{PR})} \frac{\rho_\chi(t_{PR})}{T_R^4} \frac{T_R}{\tilde{m}_B} \quad (22)$$

$$\sim 10 \times \epsilon_{CP} \epsilon_+ \frac{T_R}{g_2 \sqrt{M X_{PR}}} \quad (23)$$

$$\sim \epsilon_{CP} \epsilon_+. \quad (24)$$

Here we used $n_L(t_R)/n_L(t_{PR}) = \rho_\chi(t_R)/\rho_\chi(t_{PR})$. In reality the above estimation could be enhanced since n_L

is supplied by the gauge bosons while ρ_χ decreases more rapidly due to the decay.

Alternatively, the number density of the gauge bosons produced at the ESP can be estimated by using so-called instant preheating [17]. The simplest way discussed in Ref.[17] is to estimate the time for the instantaneous particle production $\Delta t_* \sim \chi_a / \dot{\chi}$ and use this to obtain the occupation number $n_k = \exp(-\pi k^2 / k_*^2)$, where $k_* \sim \Delta t_*^{-1}$.¹¹ Here $|\chi| < \chi_a$ denotes the region in which the adiabatic condition is violated. In the present case we find [11] $\chi_a \simeq \left(\frac{\epsilon |\dot{\chi}|^2}{\alpha g^2 M_p} \right)^{1/3}$ and

$$\begin{aligned} n_{Bs}(t_n) &= \frac{1}{2\pi^2} \int_0^\infty dk k^2 n_k \\ &\simeq \frac{k_*^3}{8\pi^3} \\ &= \frac{\lambda g^2 M_p \dot{\chi}}{8\pi^3 \xi} \sim \mathcal{O}(10^{-3}) \times g^2 M^2 \chi_n, \end{aligned} \quad (25)$$

where $n_{Bs} = 0$ is assumed just before χ passes the ESP. The above calculation is very crude but in good agreement with Eq.(89) in Ref.[11]. Now we have

$$\frac{n_+(t_R)}{T_R^3} \sim \epsilon_+ 10^{-3} \times \sqrt{\frac{M}{\chi_n}}. \quad (26)$$

Note that even in the above modest estimation the amount of the right-handed neutrinos is not negligible. This result gives a very modest estimation of the lepton asymmetry

$$\frac{n_L(t_R)}{T_R^3} \sim \epsilon_+ \epsilon_{CP} 10^{-3}, \quad (27)$$

which is still enough to realize non-thermal leptogenesis.

Before closing our discussion about leptogenesis, we have to make some comments about the hierarchy and the naturalness between the electroweak scale and M_R . The quantum correction to the electroweak scale cannot meet the naturalness criteria when M_R is as large as 10^{14} GeV. If one demands supersymmetry, the particle production played by the Higgs oscillation will be very different, since in that case there could be many directions that can break the gauge symmetry at the same time when the Higgs oscillates.

Although we calculated quantities in the simplest situation, our result is illuminating the possibility of generating baryon(lepton) number asymmetry during the preheating stage after Higgs inflation. In contrast to thermal leptogenesis the process works for $M_R \geq T_R$.

⁹ See also Ref.[19–21], in which L-R neutrino oscillation after reheating is considered for smaller M_R .

¹⁰ Here ϵ_+ is not identical to the branching ratio of the gauge boson. At the beginning of the oscillation the decay is sufficiently fast near the edge and almost all n_{Bs} is transmitted to the decay products. If we introduce the ratio ϵ_d that measures the efficiency of the decay, we can write $\epsilon_+ = \epsilon_{br} \epsilon_d$, where ϵ_{br} denotes the branching ratio. At first ϵ_+ is almost identical to the branching ratio but during the oscillation ϵ_d decreases with time.

¹¹ The power of k in the above formula n_k is correct in the standard preheating scenario, where the mass depends on the oscillating field as $m^2 \propto \chi^2$. Since in the present scenario we are considering $m^2 \propto \chi$, there could be some deviation in n_k . See Ref.[14, 15] for more details. This can change the numerical factor in the final result.

II. CONCLUSION AND DISCUSSION

Particle production is a very old idea. Many scenarios have been studied to explain experiments and observations. Among those, neutrino oscillation is a rather new idea, which is now essential for the physics in the lepton sector of the SM. Non-zero neutrino mass is needed for the neutrino oscillation, but it has to be (unnaturally) light compared with other SM particles. One way to explain the neutrino mass is to introduce the see-saw mechanism, which extends the neutrino sector by adding a missing right-handed partner of the left-handed neutrino. If the see-saw mechanism works in the neutrino sector, Higgs oscillation after Higgs inflation may lead to “another neutrino oscillation” between left and right-handed neutrinos. A strange consequence of the L-R oscillation is that the sterile neutrino (right-handed neutrino) is generated from the gauge bosons. The process is very simple. If the gauge bosons are transferred into SM-neutrinos when the mixing is significant, the SM-neutrinos can be converted into right-handed neutrinos via the left-right oscillation. The process is obviously different from the so-called off-diagonal preheating considered in past studies [12].¹² In this paper we applied the idea to the simplest model of Higgs inflation and found that leptogenesis is successful even if the reheating temperature is

¹² After integrating out the heavy modes one will find higher dimensional interaction, which has been considered for preheating

lower than the right-handed neutrino mass. At this moment our estimation is quite rough, since the particle production occurs during oscillation and important parameters are not defined local. Moreover, the process starts with the non-perturbative production of unstable gauge bosons, which may cause significant backreaction. Therefore the actual calculation is quite complicated and very careful numerical analysis is needed, especially when the flavor structure is introduced. For simplicity we used previous results in Ref.[3, 11] in our estimation and calculated the abundance of N_R from instant translation. Nevertheless, it is quite obvious that for singlet fermions the process of neutrino oscillation is much more efficient compared with the usual process of particle production [3]. Therefore, we conclude that the new mechanism we have considered in this paper is very important for the physics related to the neutrino sector.

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in Ref. [15].

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